Hadronic component of the photon spin dependent structure function g_1^γ from QCD

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Abstract

We calculate the spin dependent structure function of the polarized virtual photon $g_1^\gamma(x, Q^2, p^2)$, especially its hadronic part, using the OPE in the inverse powers of the target photon virtuality. Some model is accepted to achieve a correct analytical behaviour in the photon squared mass. This enables extrapolation to the case of a real photon.

1 Introduction

More than two decades ago it was observed [1] that the processes $e^{\pm}e^{-} \rightarrow e^{\pm}e^{-}X$ being summed over all hadronic states (see fig. 1.) allowed for the study of $\gamma\gamma$ interactions at high energies. An experimentally favourable situation for measurements of this cross section appears when one electron (corresponding to the photon with momentum q) is detected at large momentum transferred $Q^2 = -q^2$ and the other remains in a nearly forward domain $-p^2 \approx 0$. So, in practice, it is more appropriate to deal with the deep inelastic scattering (DIS) of an electron beam on a real photon target which is characterized by several nonpolarized and polarized structure functions in the same way as the usual DIS on a nucleon.

The photon structure functions in the context of the gauge theories have first been discussed by Ahmed and Ross [2]. Since 1977 [3] the deep inelastic electron scattering from a photon target has become a new subject of intensive theoretical studies in the framework of QCD. Much progress has been achieved in this direction [4]. As was firstly advocated by Witten the nonpolarized structure function $F_2(x)$ can be calculated using the perturbative QCD alone. At that time there was considerable optimism that this process was an excellent test for the perturbative QCD and might provide an accurate measurement of α_s. By now the optimism has waned considerably. This happens because Witten's suggestion is true only for the asymptotically large probe-photon momentum transfer squared where a "contact"-type term due to the photon operators in the framework of the Operator Product Expansion (OPE) turns out to be dominant. At smaller values of Q^2 the hadronic component becomes sizable, the photon admits considerable contribution that comes from the nonperturbative region. The photon has a partonic substructure induced by virtual fluctuation into the quark-antiquark pair. This process receives the short as well as long distance contribution, and the latter is due to the nonperturbative dynamics that cannot be calculated in perturbative QCD. Only if $-p^2$ is in the deep inelastic scattering region, the partonic substructure can be neglected and the photon exhibits its point-like nature. For finite $-p^2$ the target photon receives contri-

bution from the hadronic component associated with the Vector Meson Dominance (VMD) of the soft photon. Due to this contamination the hope of pure extraction of Λ_{QCD} from this reaction in a real experiment fails owing to the large uncertainty in the theoretical prediction for this part. Till recently the only estimations for the latter have been obtained from the VMD models [4].

The first QCD based calculation of the hadronic part has been initiated by Balitsky [5]. However, only a few first moments rather than x-dependence of the structure function were found in his paper. Recently a new approach for calculating the photon structure function in QCD has been developed [6]. It enables one to evaluate the structure function in the region of intermediate x and was successfully applied to the case of spin averaged scattering.

In recent times the polarized photon structure functions have attracted a lot of attention. In ref. [7] OPE and the renomalization group analysis were extended to the polarized sector, while works cited in ref. [8] deal with the first moment of the spin dependent structure function $g_1^\gamma(x)$ and its sensitivity to the chiral symmetry realization. In these papers, as in Witten's one, the hadronic component was completely disregarded. However, the knowledge of the latter is very important for comparison of the theoretical predictions with a real experiment. This holds for all presently available or foreseeable values of Q^2.

In the present study we concentrate on the calculation of the polarized structure function $g_1^\gamma(x)$ following the method mentioned above. We shall start with the consideration of the structure function when the target photon virtuality is large and spacelike $-p^2 \gg R_c^{-2}$, R_c is a confinement radius, but $-p^2 \ll Q^2$ and apply the OPE to the discontinuity [9] of the forward photon-photon scattering amplitude which results in the expansion in the inverse powers of target photon virtuality. On the other hand, using the analytical properties of the structure function in the photon off-shellness $-p^2$ it is represented via the dispersion relation through the contribution of the vector meson and continuum. Comparing the two representations of the same quantity we can fix unambiguously all parameters of hadronic spectrum.

The structure function obtained possesses correct analytical properties in the target photon mass and has no fictitious kinematical singularities inherent in the perturbative diagrams.

The sensitivity to the QCD radiative corrections is poor until very large Q^2 is attained. Thus, we restrict ourselves to the lowest order graphs and to the consideration of light quarks only; the result is correct for intermediate $Q^2 \leq 10GeV^2$. However, for too low Q^2 the higher twist corrections may be important.

2 OPE for the four-point amplitude

To start with we consider the four-point correlation function

$$T_{\mu\nu,\alpha\beta} = 4\pi\alpha_{em} i^3 \int d^4x d^4y d^4z e^{iqx+ip(y-z)} \langle 0|T\{j_\mu(x) j_\nu(0) j_\alpha(y) j_\beta(z)\}|0\rangle, \qquad (1)$$

where $j_\mu = \sum_q Q_q \bar{\psi}_q \gamma_\mu \psi_q$ is the electromagnetic quark current. This amplitude is originated from the T-product of two electromagnetic currents between the photon states and application of the Lehmann-Symanzik-Zimmermann reduction formula. The discontinuity across the branch cut on the real axis in the complex plane of $\omega = \frac{1}{x}$, where x is the usual Bjorken variable, gives us the structure function we are interested in.

Note that due to the weak coupling of the photon to the hadronic states there is no need to extract the physical state of interest with the help of some auxiliary procedure, *e.g.* the Borel transformation.

In order to find the polarized spin structure function, we isolate it as a coefficient in front of an appropriate tensor structure, namely

$$\frac{1}{\pi} Im T_{\mu\nu,\alpha\beta} \epsilon_\alpha \epsilon^*_\beta = \frac{i}{(pq)} \epsilon_{\mu\nu\lambda\sigma} q_\lambda s_\sigma g_1^\gamma(x, Q^2, p^2). \qquad (2)$$

where $s_\sigma = i\epsilon_{\alpha\beta\gamma\sigma} \epsilon_\alpha \epsilon^*_\beta p_\gamma$ and ϵ_α is a photon polarization vector. More precisely, for the purposes of this paper it is enough to pick out the antisymmetric tensor $(g_{\mu\alpha} g_{\nu\beta} - g_{\mu\beta} g_{\nu\alpha})$.

As a first step we have to find the contribution of the unit operator. This result is well known since the "photon-photon fusion" process was calculated even before the advent of QCD [10]. However, to reproduce unambiguously the spectral densities in the dispersion representation for the structure function, which require some intermediate result, such a calculation has to be performed over again. We restrict ourselves to the scaling approximation, i.e. to taking into account the first nonvanishing term in the expansion in powers of p^2/Q^2, therefore limiting to the leading twist-2 contribution. The result is

$$
\begin{aligned}
&g_1^{\gamma}(x,Q^2,p^2)_{pert} \\
&\quad= \frac{\alpha_{em}}{\pi} N_c N_f \langle Q_q^4\rangle \left\{ \left[2-3x+[x^2-(1-x)^2]\int_0^{Q^2/x^2} \frac{dp'^2 p'^2}{(p'^2-p^2)^2} \right] - (1-x) \right\} \\
&\quad= \frac{\alpha_{em}}{\pi} N_c N_f \langle Q_q^4\rangle [x^2-(1-x)^2] \left[\ln\left(\frac{Q^2}{-p^2x^2}\right) - 2 \right], \qquad (3)
\end{aligned}
$$

where $\langle Q_q^4\rangle = \frac{1}{N_f}\sum_q Q_q^4$ is an average of the fourth powers of the quark charges. This result is twice that represented by diagrams in fig. 2 due to the clockwise and counter-clockwise directions of the internal quark lines, each term in the curly brackets corresponding to graph (a) and (b), respectively. The first line of this equation will be used in the following to fix the parameters of the hadronic spectrum.

From all power corrections up to dimension eight we calculate only one due to the gluon condensate $\langle \frac{\alpha_s}{\pi} G^2\rangle$. This can be elucidated by the facts that the lowest dimension quark condensate $\langle\bar\psi\psi\rangle$ cannot appear due to chiral invariance as it is accompanied by the light quark mass which we set equal to zero in all calculations. The contribution of the three-gluon condensate $\langle g^3 f G^3\rangle$ is usually small. The dimension six four-quark condensate can be omitted because its contribution is proportional to the delta function — $\delta(1-x)$ and turns out to be beyond the scope of the method.

To simplify the calculation of the leading power correction, it is convenient to use the fixed-point gauge for the background gluon field $(x-x_0)_\mu B_\mu^a(x)=0$. We chose the fixed point in the vertex of the hard photon emission $x_0=0$. The quark

propagator in this gauge up to the order $O(G^2)$ looks like [11]

$$S(x,y) = \int \frac{d^4k}{(2\pi)^4} e^{-ik(x-y)} \Big\{ \frac{\not k}{k^2} + g\widetilde{G}^a_{\alpha\beta} t^a \frac{k_\alpha}{k^4}\gamma_\beta\gamma_5 + \frac{1}{2} g G^a_{\alpha\beta} t^a y_\alpha \left[\frac{\gamma_\beta}{k^2} - 2\frac{k_\beta \not k}{k^4}\right] - \frac{\langle g^2G^2\rangle}{3^2 2^5}\left[\frac{y^2 \not k}{k^4} + 4\frac{\not k (ky)^2}{k^6} - 2\frac{\not y (ky)}{k^4}\right]\Big\}. \qquad (4)$$

A non-zero contribution comes in the leading twist from the diagrams depicted in fig. 3 and each term in the curly brackets corresponds to the diagrams (a), (b) and (c), respectively:

$$\begin{aligned} & g_1^\gamma(x,Q^2,p^2)_{\langle \frac{\alpha_s}{\pi}G^2\rangle} \\ &= \frac{\alpha_{em}}{\pi} N_c N_f \langle Q_q^4\rangle \frac{1}{36}\frac{\pi^2}{p^4}\langle \frac{\alpha_s}{\pi}G^2\rangle \left\{\left[\frac{8}{3}\frac{1}{x^2} - \frac{4}{3}\frac{1}{x}\right] + \left[\frac{8}{3}\frac{1}{x^2}\right] + \left[\frac{4}{3}\frac{1}{x}\right]\right\} \\ &= \frac{\alpha_{em}}{\pi} N_c N_f \langle Q_q^4\rangle \frac{4}{27}\frac{\pi^2}{p^4}\langle \frac{\alpha_s}{\pi}G^2\rangle \frac{1}{x^2}. \end{aligned} \qquad (5)$$

Collecting all contributions we obtain the following structure function for the off-shell polarized target photon:

$$\begin{aligned} & g_1^\gamma(x,Q^2,p^2) \\ &= \frac{\alpha_{em}}{\pi} N_c N_f \langle Q_q^4\rangle \left\{[x^2 - (1-x)^2]\left[\ln\left(\frac{Q^2}{-p^2x^2}\right) - 2\right] + \frac{4}{27}\frac{\pi^2}{p^4}\langle \frac{\alpha_s}{\pi}G^2\rangle \frac{1}{x^2}\right\}. \end{aligned} \qquad (6)$$

The singularity in the gluon condensate contribution shows invalidity of the method in the region of small x. The applicability region of this formula, which accounts for the hadronic part, can be found by usual requirement that the power corrections would not exceed 50% of the main perturbative term. Going as lower as to the point $-p^2 = m_\rho^2 = 0.6 GeV^2$ (ρ-meson is the lowest prominent state in the channel with the photon quantum numbers) in eq. (6) the gluon condensate comprises less than 50% only for $x \geq 0.6$. The upper limit of permitted x-values will be estimated below. From these facts it follows that the present approach could not be used for the calculation of the moments of the structure function.

Now we are in a position to make some comments on the nonlocal condensates.

Introducing the nonlocal gluon condensate it is probably impossible to get rid of singular behaviour of the corresponding contribution to the structure function [12]; the former diminishes the degree of singularity by one power in the $x \sim 0$ region.

However, the singular δ-type behaviour of the quark condensate contribution can be smeared over the whole region of the momentum fraction from zero to unity by introducing the concept of a nonlocal quark condensate. Such an attempt was made in ref. [13] for the spin averaged structure function $F_2(x)$. But there are two shortcomings in this paper. First, the authors claim that the contribution due to the quark condensate improves considerably the description of experimental data, though, there is a numerical error in their answer: the coefficient in front of the nonlocal vector quark condensate is three times smaller. The second fact is connected with improper treatment of nonlocal objects. The diagram with a nonlocal scalar condensate used in their paper does not exist. The arguments are as follows. The philosophy of QCD sum rules is the mutual relationship between the resonance and asymptotic parameters. The leading asymptotic behaviour is given by perturbative QCD while nonperturbative contributions enter through the power corrections which die out for asymptotically large momenta squared of the problem — in the formal limit $-p^2, -\tau^2 \to \infty$, where we introduce nonzero t-channel momentum τ in order to treat the factorization of small and large distances properly. The technique is to extract the coefficient function that can be dealt perturbatively — the part of the diagram with hard momentum flow, while the soft lines are parametrized via vacuum condensates that accumulate nonperturbative information. In the diagram in ref. [13] there is no short-distance coefficient function. This can be traced by the fact that in the local limit it becomes disconnected. To preserve the latter property it is necessary to attach the gluon line connecting the upper and lower parts (see fig. 4(a)). However, in the small $-\tau^2$ kinematics, actually, we are interested in the case of $-\tau^2 = 0$, the diagram will behave as $\frac{1}{-\tau^2}$ for the local quark condensates or something like $\frac{1}{-\tau^2+\lambda_q^2}$, where $\frac{1}{\lambda_q}$ is a correlation length of quarks in the vacuum, for the nonlocal one. For small $-\tau^2$ the use of the bare gluon propagator is no longer reliable as the average virtuality $\lambda_q^2 = 0.4GeV^2$ [14] and turns outside the perturbative domain. We expect that for low $-\tau^2$ the nonperturbative effect does provide an appropriate infrared (IR) cut-off: the perturbative behaviour

$\frac{1}{-\tau^2}$ is substituted by the mass of an appropriate meson $\frac{1}{m_R^2}$, that is the "natural" scale for the boundary between the non- and perturbative regions. In principle, this program could be traced by reformation of the original OPE [5] to the case when the momentum $-\tau^2$ can be arbitrary small (even zero). However, there is an essential difficulty as compared to the form factor problem [15], namely, due to the presence of the t-channel non-analyticities or singularities in each moment of the structure function, we need an infinite number of parameters to be found from additional sum rules which enter the model of the bilocal power corrections. Obviously, this is an impossible task. We will not go into details here referring the interested reader to the paper [16] where this question was studied for the case of quark distribution in the pion. Therefore, limiting ourselves to the central region in the Bjorken variable we discard these terms completely. Of course, there are other diagrams (see fig. 4(b)) which are regular in $-\tau^2$ but we believe that the contribution neglected could not affect seriously the final result.

3 Polarized photon structure function

We can use the analytical properties in p^2 [17] and represent the structure function via a dispersion relation with respect to p^2 in terms of physical states

$$g_1^\gamma(x,Q^2,p^2) = G_0(x) + \int_0^\infty dp'^2 \frac{G_1(x,p'^2)}{(p'^2-p^2)} + \int_0^\infty dp_1'^2 \int_0^\infty dp_2'^2 \frac{G_2(x,p_1'^2,p_2'^2)}{(p_1'^2-p^2)(p_2'^2-p^2)}. \tag{7}$$

To calculate the functions G_i, we use a standard technique in QCD sum rules, viz., the simple model of a lowest resonance plus continuum

$$\begin{aligned} G_1(x,p'^2) &= G_1^{(1)}(x)\delta(p'^2-m_\rho^2) + G_1^{(2)}(x)\theta(p'^2-s_0), \\ G_2(x,p_1'^2,p_2'^2) & \\ &= G_2^{(1)}(x)\delta(p_1'^2-m_\rho^2)\delta(p_2'^2-m_\rho^2) + G_2^{(2)}(x)\theta(p_1'^2-s_0)\theta(p_2'^2-s_0), \end{aligned} \tag{8}$$

where $s_0 = 1.5GeV^2$ is a threshold value for the vector meson channel and $m_\rho^2 = 0.6GeV^2$ is a ρ-meson mass squared. In other words, we provide a natural cut-off

on the transverse momentum in the loop $\theta(k_\perp^2 - x(1-x)s_0^2)$ attributing the region of small $k_\perp^2$ to the nonperturbative contribution. The continuum threshold s_0 has been found from the analysis of two-point correlators for ρ-meson in the framework of the QCD sum rules [18].

Requiring that at $-p^2 \to \infty$ eq.(7) must coincide with the bare quark loop, we obtain:

$$
\begin{aligned}
&G_0(x) = -\frac{\alpha_{em}}{\pi} N_c N_f \langle Q_q^4 \rangle [x^2 - (1-x)^2], \\
&G_1^{(2)}(x) = 0, \\
&G_2^{(2)}(x) = \frac{\alpha_{em}}{\pi} N_c N_f \langle Q_q^4 \rangle [x^2 - (1-x)^2] {p_1'}^2 \delta({p_1'}^2 - {p_2'}^2) \theta(Q^2/x^2 - {p_1'}^2)
\end{aligned}
\tag{9}
$$

Substituting them back into the dispersion relation (7) and expanding in the inverse powers of p^2 we can compare the resulting expression with the QCD calculated $g_1^\gamma(x, Q^2, p^2)$ (6) and fix the remaining unknown functions $G_1^{(1)}(x)$ and $G_2^{(1)}(x)$, namely:

$$
\begin{aligned}
&G_1^{(1)}(x) = 0, \\
&G_2^{(1)}(x) = \frac{\alpha_{em}}{\pi} N_c N_f \langle Q_q^4 \rangle \left[[x^2 - (1-x)^2] \frac{s_0^2}{2} + \frac{4}{27} \pi^2 \langle \frac{\alpha_s}{\pi} G^2 \rangle \frac{1}{x^2} \right].
\end{aligned}
\tag{10}
$$

Finally, we collect all functions and perform the momentum integration in eq.(7) keeping only the twist-2 contribution. We come to the polarized virtual structure function which possesses the correct analytical properties in the photon squared mass and accounts for the hadronic part:

$$
\begin{aligned}
&g_1^\gamma(x, Q^2, p^2) = \frac{\alpha_{em}}{\pi} N_c N_f \langle Q_q^4 \rangle \\
&\Bigg\{ -[x^2 - (1-x)^2] + [x^2 - (1-x)^2] \left[\ln \left(\frac{Q^2}{x^2 (s_0 - p^2)} \right) + \frac{p^2}{(s_0 - p^2)} \right] \\
&+ \frac{1}{(p^2 - m_\rho^2)^2} \left[[x^2 - (1-x)^2] \frac{s_0^2}{2} + \frac{4}{27} \pi^2 \langle \frac{\alpha_s}{\pi} G^2 \rangle \frac{1}{x^2} \right] \Bigg\}.
\end{aligned}
\tag{11}
$$

4 Result and conclusion

As we noticed above, the correction due to the gluon condensate comprises no more than 50% only for $x \geq 0.6$ at $-p^2 = m_\rho^2$. Therefore, for this x-values we can ex-

trapolate the polarized virtual structure function given by eq. (11) to the point $p^2 = 0$. In the large-x region, we could not trust our model for hadronic spectrum. As was shown in ref. [6] as $x \to 1$ multihadron states exceed the contribution of single ρ-meson. Moreover, since the function $G_2^{(1)}(x)$ given by eq. (10) is related to the polarized structure function of the meson[1] $\Delta f_\rho(x)$ it should reproduce $(1-x)$-behaviour as $x \to 1$ governed by the quark counting rules. However, it is easy to check that it does not take place. The upper limit of accepted x-values may be estimated from he requirement that the created hadron state should exceed considerably the ρ-meson mass $M_X^2 \gg m_\rho^2$. Then from the kinematical constraint $x = 1/(1+(M_X^2 - m_\rho^2)/Q^2)$ we have the restriction $x < 0.8$. Therefore, the approach for calculating the polarized photon structure function $g_1^\gamma(x)$ is valid in the interval $0.6 \leq x < 0.8$. This situation is quite similar to that for calculating the polarized proton structure functions in the QCD sum rules framework where the range of permitted x-values is also very narrow [19].

Note that from the point of view of the QCD sum rules an interpretation of separate contributions to the structure function is different from the naïve division of the latter in terms of the perturbative (found from the box diagrams) and VMD parts which inevitably contains double counting. In the language of the present approach the part of the perturbative loop is dual to the lowest resonance with corresponding quantum numbers and duality interval s_0, while the remaining part of the loop corresponds to the higher states contribution (the continuum). The resonance also receives the contribution from the nonperturbative power corrections.

In fig. 5 we represent the result of calculation of the real photon polarized structure function. The long- and short-dashed curves correspond to the continuum and the vector-meson contributions to the latter, respectively. They are represented by the first and second terms in the expression (11). The solid curve is a full structure function (11). We have used the standard value for the gluon condensate $\langle \frac{\alpha_s}{\pi} G^2 \rangle =$

[1] Eq. (10) corresponds to the local duality relation for the structure function of the polarized ρ-meson.

$0.012 GeV^4$ [18]. From fig. 5. it is clearly seen that the hadronic component obtained is very large.

Unfortunately, up to now there are no available experimental data on the photon spin structure function g_1^γ as its measurement is at the limits of the possibilities of the polarized colliders. Future measurements would provide an important insight into the underlying dynamical effects associated with the polarized quark and gluon content of the photon.

Acknowledgments

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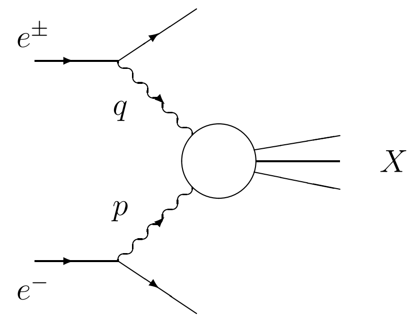

Figure 1: The kinematics of the two-photon reaction $e^{\pm}e^{-} \to e^{\pm}e^{-}X$.

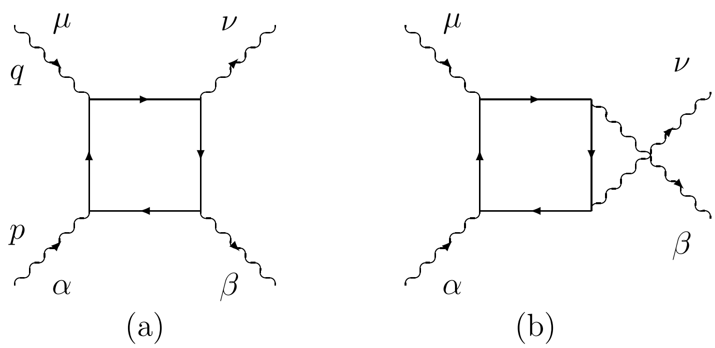

Figure 2: Perturbative diagrams for the unit operator in the operator expansion.

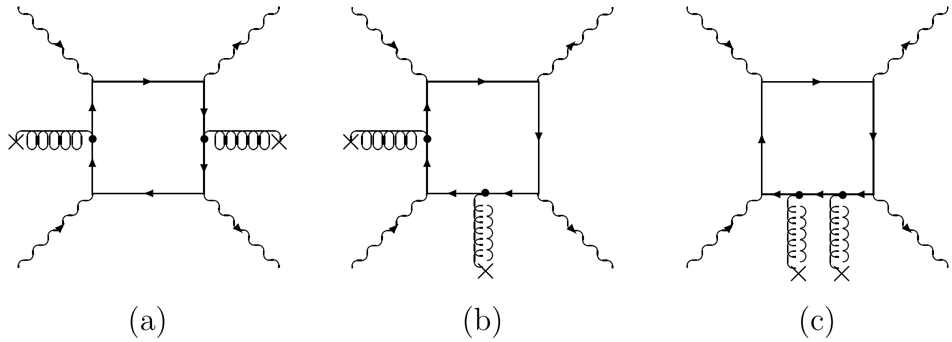

Figure 3: Gluon condensate contribution to the imaginary part of the forward $\gamma\gamma$-scattering amplitude.

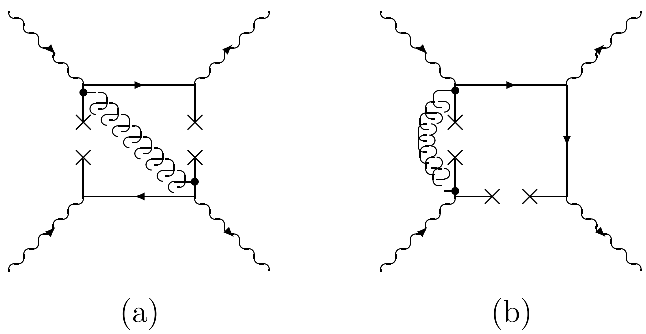

Figure 4: Four-quark condensate contribution to the imaginary part of the non-forward $\gamma\gamma$-scattering amplitude.

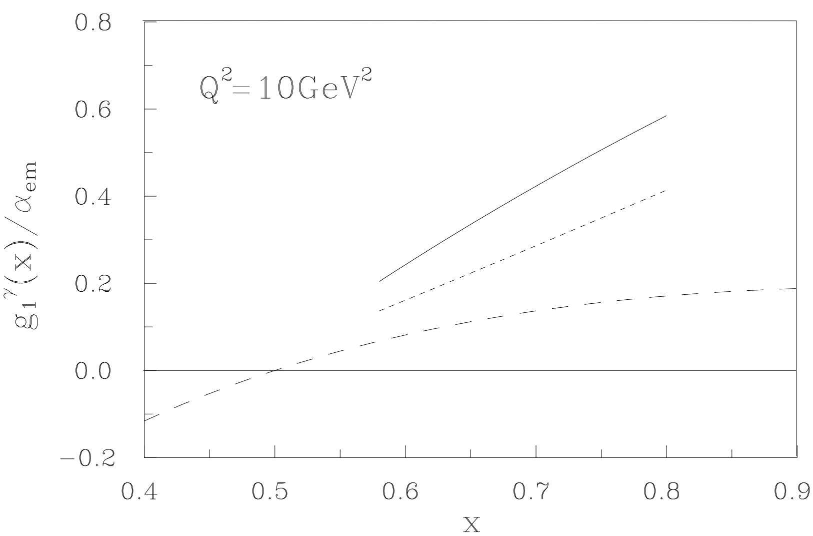

Fig. 5. Spin dependent structure function of the real photon at $Q^2 = 10GeV^2$. The solid curve corresponds to the full structure function given by eq. (11), while long- and short-dashed lines correspond to the continuum and hadronic contributions to the latter.